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CAN GLUONS TRACE BARYON NUMBER ?

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Abstract

QCD as a gauge non-Abelian theory imposes severe constraints on the structure of the baryon wave function. We point out that, contrary to a widely accepted belief, the traces of baryon number in a high-energy process can reside in a non-perturbative configuration of gluon fields, rather than in the valence quarks. We argue that this conjecture can be tested experimentally, since it can lead to substantial baryon asymmetry in the central rapidity region of ultra-relativistic nucleus-nucleus collisions.

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In QCD, quarks carry colour, flavour, electric charge and isospin. It seems only natural to assume that they also trace baryon number. However, this latter assumption is not dictated by the structure of QCD, and therefore does not need to be true. Indeed, the assignment of the baryon number $B = 1/3$ to quarks is based merely on the naive quark model classification. But any physical hadron state in QCD should be represented by a state vector which is gauge-invariant – the constraint which is ignored in most of the naive quark model formulations. This constraint turns out to be very severe; in fact, there is only one way to construct a gauge-invariant state vector of a baryon from quarks and gluons [1] (note however that there is a large amount of freedom in choosing the paths connecting x to x_i):

$$B = \epsilon^{ijk} \left[Pexp \left(ig \int_{x_1}^x A_\mu dx^\mu \right) q(x_1) \right]_i \left[Pexp \left(ig \int_{x_2}^x A_\mu dx^\mu \right) q(x_2) \right]_j \times \left[Pexp \left(ig \int_{x_3}^x A_\mu dx^\mu \right) q(x_3) \right]_k. \quad (1)$$

The “string operators” in (1) acting on the quark field $q(x_n)$ make it transform as a quark field at point x instead of at x_n . The ϵ tensor then constructs a local colour singlet and gauge invariant state out of three quark fields (see Fig.1a). The B in eq. (1) is a set of gauge invariant operators representing a baryon in QCD. With properly optimised parameters it is used extensively in the first principle computations with lattice Monte Carlo attempting to determine the nucleon mass. The purpose of this work is to study its phenomenological impact on baryon number production in the central region of nucleus-nucleus collisions.

It is evident from the structure of (1) that the trace of baryon number should be associated not with the valence quarks, but with a non-perturbative configuration of gluon fields located at the point x - the “string junction” [1]. This can be nicely illustrated in the string picture: let us pull all of the quarks away from the string junction, which we keep fixed at point x . This will lead to $\bar{q}q$ pair production and string break-up, but the baryon will always restore itself around the string junction. The quark composition of this resulting baryon will in general differ from the composition of the initial baryon. It is important to note that the assignment of the trace of baryon number to gluons is not only a feature of a particular kind of the string model, but is a consequence of the local gauge invariance principle applied to baryons.

Surprisingly, at first glance this non-trivial structure of baryons in QCD does not seem to reveal itself in hadronic interactions in any appreciable way. This is because to probe the internal structure of baryons, we need hard interactions, which usually act on single quarks. The baryon number therefore can always be associated with the remaining diquark, without specifying its internal structure. To really understand what traces the baryon number, one needs therefore to

study processes in which the flow of baryon number can be separated from the flow of valence quarks.

In this Letter we suggest that studies of baryon production in the central rapidity region of ultra-relativistic pp and especially AA collisions provide a crucial possibility to test the baryon structure. Our findings may prove important for understanding the properties of dense QCD matter produced in AA collisions at high energies. In what follows, we shall first formulate our ideas qualitatively, and then give some quantitative estimates based on the topological expansion of QCD [1, 2] and Regge phenomenology.

Let us consider first an ultra-relativistic pp collision in its centre-of-mass frame, which coincides with the lab frame in collider experiments. At sufficiently high energies, the valence quark distributions will be Lorentz-contracted to thin pancakes with the thickness of

$$z_V \simeq \frac{1}{x_V P}, \quad (2)$$

where P is the c.m. momentum in the collision, and $x_V \sim 1/3$ is a typical fraction of the proton's momentum carried by a valence quark. The typical time needed for the interaction of valence quarks from different protons with each other during the collision is given by the characteristic interquark distance in the impact parameter plane, $t_{int} = \text{const} \sim O(1 \text{ fm})$. However, the time available for this interaction in the collision is only $t_{coll} \sim z_V \sim (x_V P)^{-1}$. It is therefore clear that at sufficiently high energies, when $t_{coll} \ll t_{int}$, the valence quarks of the colliding protons do not have time to interact during the collision and go through each other, populating the fragmentation regions. In the conventional picture, the baryon number follows the valence quarks.

At first glance, the argument looks correct, and is well supported experimentally - the leading effect for baryons in high energy pp collisions is well established. However, the structure of the gauge-invariant baryon wave function (1) suggests that this scenario may not be entirely consistent. As we have stressed above, in QCD the trace of the baryon number has to be associated with the non-perturbative configuration of the gauge field. This “string junction” contains an infinite number of gluons which, therefore, by virtue of momentum conservation, should carry on the average an infinitely small fraction $x_S \ll x_V$ of the proton's momentum. We therefore expect that the “string junction” configuration may not be Lorentz-contracted to a thin pancake even at asymptotically high energies (see Fig.2), since

$$z_S \simeq \frac{1}{x_S P} \gg z_V. \quad (3)$$

In this case the string junction will always have enough time to interact, and we may expect to find stopped baryons in the central rapidity region even in a high-

energy collision (of course there will also be baryon-antibaryon pair production in addition). This argument leads to a peculiar picture of a high-energy pp collision: in some events, one or both of the string junctions are stopped in the central rapidity region, whereas the valence quarks are stripped-off and produce three-jet events in the fragmentation regions. Immediately after collision, the central region is then filled by a gluon sea containing one or two twists, which will later on be dressed up by sea quarks and will form baryon(s). Note that the quark composition of the produced baryons will in general differ from the composition of colliding protons.

Why then is the leading baryon effect a gross feature of high-energy pp collisions? The reason may be the following. The string junction, connected to all three of the valence quarks, is confined inside the baryon, whereas pp collisions become on the average more and more peripheral at high energies. Therefore, in a typical high-energy collision, the string junctions of the colliding baryons pass far away from each other in the impact parameter plane and do not interact. One can however select only central events, triggering on high multiplicity of the produced hadrons. In this case, we expect that the string junctions will interact and may be stopped in the central rapidity region. This should lead to the baryon asymmetry in the central rapidity region: even at very high energies, there should be more baryons than antibaryons there.

Fortunately, the data needed to test this conjecture already exist: the experimental study of baryon and antibaryon production with trigger on associated hadron multiplicity has been already performed at ISR, at the highest energy ever available in pp collisions [3]. This study has revealed that in the central rapidity region, the multiplicities associated with a proton are higher than with an antiproton by $\simeq 10\%$. It was also found that the number of baryons in the central rapidity region substantially exceeds the number of antibaryons [4]. These two observations combined indicate the existence of an appreciable baryon stopping in central pp collisions even at very high energies [3].

Where else do we encounter central baryon-baryon collisions? In a high energy nucleus-nucleus collision, the baryons in each of the colliding nuclei are densely packed in the impact parameter plane, with an average inter-baryon distance

$$r \simeq \rho^{-1/2} A^{-1/6}, \quad (4)$$

where ρ is the nuclear density, and A is the atomic number. The impact parameter b in an individual baryon-baryon interaction in the nucleus-nucleus collision is therefore effectively cut off by the packing parameter: $b \leq r$. In the case of a lead nucleus, for example, r appears to be very small: $r \simeq 0.4 \text{ fm}$, and a central lead-lead collision should therefore be accompanied by a large number of interactions among the string junctions. This may lead to substantial baryon stopping even at RHIC and LHC energies.

We shall now proceed to more quantitative considerations. In the topological expansion scheme [1], the separation of the baryon number flow from the flow of valence quarks in baryon-(anti)baryon interaction can be represented through a t -channel exchange of the quarkless junction-antijunction state with the wave function given by

$$M_0^J = \epsilon_{ijk} \epsilon^{i'j'k'} \left[Pexp \left(ig \int_{x_1}^{x_2} A_\mu dx^\mu \right) \right]_{i'}^i \left[Pexp \left(ig \int_{x_1}^{x_2} A_\mu dx^\mu \right) \right]_{j'}^j \times \left[Pexp \left(ig \int_{x_1}^{x_2} A_\mu dx^\mu \right) \right]_{k'}^k. \quad (5)$$

The structure of the wave function (5) is illustrated in Fig.1b - it is a quarkless closed string configuration composed from a junction and an antijunction. In the topological expansion scheme, the states (5) lie on a Regge trajectory; its intercept can be related to the baryon and reggeon intercepts [1]:

$$\alpha_0^J(0) \simeq 2\alpha_B(0) - 1 + 3(1 - \alpha_R(0)) \simeq \frac{1}{2}, \quad (6)$$

where the baryon intercept $\alpha_B(0)$ has been set equal to 0, and the reggeon intercept $\alpha_R(0)$ equal to $1/2$.

The M_0^J exchange should dominate the proton-antiproton annihilation at high energies [1]. Indeed, annihilation requires the baryon number transfer in the t -channel. Conventionally, this corresponds to the baryon exchange with the intercept $\alpha_B(0) \simeq 0$. Since in Regge theory the energy dependence of the cross section is given by $s^{\alpha(0)-1}$, and $\alpha_0^J(0) > \alpha_B(0)$, the M_0^J exchange should give the dominant contribution at high energies (see Fig.3), leading to the following energy dependence of the annihilation cross section:

$$\sigma_{\bar{p}p}^{ann} \sim \left(\frac{s}{s_0} \right)^{\alpha_0^J(0)-1} \simeq \left(\frac{s}{s_0} \right)^{-1/2} \quad (7)$$

instead of the s^{-1} dependence implied by conventional baryon exchange ($s_0 \simeq 1 \text{ GeV}^2$ is the usual parameter of Regge theory). Up to the highest energies where the annihilation can still be experimentally distinguished from other inelastic processes, the energy dependence (7) is confirmed by the data. Moreover, the entire difference between the total pp and $\bar{p}p$ cross sections can be attributed to annihilation (see [5] for a recent review).

Let us now turn to the consideration of baryon stopping in pp collisions. The relevant diagrams are shown in Fig.4; we consider the simultaneous stopping of the two string junctions in the central rapidity region, accompanied by three-jet events in the fragmentation regions (Fig.4a), and the stopping of the junction of one proton in the soft parton field of the other, accompanied by one three-jet

event (Figs.4b,c). To calculate the cross sections, it is convenient to evaluate the discontinuity of the corresponding three-particle elastic scattering process [6, 7] (see Fig.5). Introducing the four-momentum p_B of the produced baryon (or the total momentum of the pair of baryons in the diagram of Fig.5a) and the four momenta of the colliding protons p_1, p_2 we have the following expressions for the invariant energy in the proton-baryon systems:

$$\begin{aligned} s_1 &\simeq \sqrt{s} m_t e^{-y^*}, \\ s_2 &\simeq \sqrt{s} m_t e^{y^*}, \end{aligned} \quad (8)$$

where $s = (p_1 + p_2)^2$ is the c.m.s. energy squared of the pp collision, y^* is the c.m.s. rapidity of the produced baryon(s), and m_t is its transverse mass $m_t^2 = m_B^2 + p_B^2$. The product of the invariants (8) satisfies the relation

$$s_1 s_2 \simeq m_t^2 s. \quad (9)$$

Let us denote the coupling of the M_0^J reggeon and pomeron to the proton by G_p^M and G_p^P respectively, and introduce scalar functions $f_B^{MM}(m_t^2)$ and $f_B^{MP}(m_t^2)$ describing the “two baryons - M_0^J - M_0^J ” and “one baryon - M_0^J - Pomeron” vertices. The standard calculation [6, 7] then allows us to calculate the cross sections; for the diagram of Fig.5a we get

$$E_B \frac{d^3 \sigma^{(2)}}{d^3 p_B} = 8\pi [G_p^M(0)]^2 f_B^{MM}(m_t^2) \left(\frac{\sqrt{s} m_t}{s_0} \right)^{2\alpha_0^J(0)-2}. \quad (10)$$

Analogous calculation for the sum of diagrams in Figs.5b,c gives

$$\begin{aligned} E_B \frac{d^3 \sigma^{(1)}}{d^3 p_B} &= 8\pi G_p^M(0) G_p^P(0) f_B^{MP}(m_t^2) \left(\frac{\sqrt{s} m_t}{s_0} \right)^{\alpha_0^J(0) + \alpha_P(0) - 2} \\ &\times \left(\exp[y^*(\alpha_P(0) - \alpha_0^J(0))] + \exp[-y^*(\alpha_P(0) - \alpha_0^J(0))] \right). \end{aligned} \quad (11)$$

Using the value (6) of the M_0^J intercept, we find that the double baryon production cross section (10) has $\sim s^{-1/2}$ energy dependence and (within the central rapidity region, where our considerations apply) does not depend on rapidity. The cross section (11) of single baryon stopping also decreases with energy, but much more slowly. Writing down the Pomeron intercept as

$$\alpha_P(0) = 1 + \Delta, \quad (12)$$

one gets the energy dependence of the cross section (11) in the form $\sim s^{-1/4 + \Delta/2}$. At very high energies the process of single baryon stopping will therefore be more important. Note however that the ISR data show a large value of the correlation between the probabilities of the stopping of the beam baryons, indicating that

the process of double baryon stopping may still dominate even at rather high energies. The rapidity dependence of (11) does not show a “central plateau”, indicating instead a “central valley” structure. This is in accord with the ISR data; moreover the particular rapidity dependence of (11) reproduces the data [8] reasonably well.

The functional dependence of the single baryon stopping cross section similar to (11) has been advocated before [9] in a different approach; the authors considered a specific mechanism of perturbative destruction of the fast diquark, accompanied by the baryon number flow over a large rapidity gap. In the framework of their approach, the authors of ref.[9] have also performed a calculation of the cross section, based on the combination of perturbative technique and constituent quark model.

Unfortunately, since we believe that the dynamics of the stopping process is genuinely non-perturbative, so far we have not been able to find a reliable way of computing the constant $G_p^M(0)$ entering the expression (11) and can only extract it from the existing data. However once it is done, we can perform parameter-free extrapolation to higher energies.

As we have already stressed above, the formulae (10,11) refer to the *net* baryon number, i.e. they refer to the difference between the number of produced baryons and antibaryons. The process of baryon-antibaryon pair production will therefore represent an important background to baryon stopping; the data [8] show that at ISR energies the probability of baryon stopping is about three times smaller than the probability of baryon pair production. At high energies, the dominant contribution to the $\bar{B}B$ pair production will be given by the interaction of two Pomerons, with the cross section

$$E_B \frac{d^3\sigma^{(\bar{B}B)}}{d^3p_B} = 8\pi[G_p^P(0)]^2 f_{\bar{B}B}^{PP}(m_t^2) \left(\frac{\sqrt{s} m_t}{s_0} \right)^{2\alpha_P(0)-2}, \quad (13)$$

representing an energy-independent fraction of the total cross section:

$$\sigma^{\bar{B}B} \sim \sigma^{tot} \sim \left(\frac{s}{s_0} \right)^{\alpha_P(0)-1}. \quad (14)$$

Since the topological structure of the Pomeron is that of a cylinder [1]

$$P = Tr \left[\left(P \exp \left(ig \oint A_\mu dx^\mu \right) \right) \right], \quad (15)$$

at high energies the associated multiplicities in the processes of single ($n^{(1)}$) and double ($n^{(2)}$) baryon stopping will be higher than in the average inelastic event, described by the cut of the Pomeron exchange diagram (see Fig.6):

$$n^{(1)} \simeq \frac{5}{4} n^{inel}, \quad n^{(2)} \simeq \frac{3}{2} n^{inel}, \quad (16)$$

as we already discussed above at the qualitative level. Also, since the baryon stopping and baryon pair production arise in our scheme from different kinds of t -channel exchange (M_0^J and the Pomeron, respectively), we expect a small correlation between the events with the pair production and stopping. Experimentally, it was found to be 0.16 ± 0.22 [8].

Before we turn to the discussion of nucleus-nucleus collisions, it is useful to recall the geometrical picture of high energy scattering in the impact parameter plane. In the impact parameter representation, the growth of the total cross section at high energies as described by one-Pomeron exchange can be attributed to the increase of the effective radius of the interaction, according to

$$R_{int} \simeq \sqrt{2R_p^2 + \alpha'_P \ln(s/s_0)}, \quad (17)$$

where R_p is a constant and α'_P is the slope of the Pomeron trajectory. (The Froissart bound allows even faster growth of the interaction radius with energy: $R_{int} \sim \ln(s/s_0)$.) The central region of the disk becomes completely black at high energy if $\alpha_P(0) > 1$; a further growth of the interaction strength in the centre of the disk is prevented by the unitarity constraint imposed on the partial amplitudes.

The colliding nuclei, the transverse plane nucleon distributions in which are characterized by the packing parameter (4) will therefore see each other as uniform black disks. This means that in a *central* nucleus-nucleus collision the cross section of the inelastic nucleon-nucleon collision will not further increase with energy when $R_{int} \gg r$ since the soft peripheral interactions building up the Pomeron will be effectively screened out. On the other hand, the processes of baryon stopping are central in the impact parameter plane, and therefore may not be screened in the case of nuclear collisions. A very slow decrease of the cross section (11) with energy implies then that even at LHC energies the nuclear stopping may still be present, as we shall now discuss.

The ISR data [10, 11] show that at $\sqrt{s} = 53 \text{ GeV}$ the cross sections of proton and antiproton production at $y^* = 0$ and $p_t = 0.6 \text{ GeV}/c$ are

$$\frac{d^3\sigma^p}{d^3p}(y^* = 0) = 0.700 \pm 0.162 \text{ mb GeV}^{-2}; \quad \frac{d^3\sigma^{\bar{p}}}{d^3p}(y^* = 0) = 0.430 \pm 0.033 \text{ mb GeV}^{-2}. \quad (18)$$

Since the mechanism of baryon-antibaryon pair production (see (13)) obviously leads to an equal number of protons and antiprotons, the data (18) imply that the following fraction of the protons in the central rapidity region is produced by stopping:

$$f_{st}(\sqrt{s} = 53 \text{ GeV}) = \frac{\sigma^p - \sigma^{\bar{p}}}{\sigma^p + \sigma^{\bar{p}}} \simeq 25\%. \quad (19)$$

We can now estimate the ratio R of multiplicities associated with proton (n^p) and antiproton ($n^{\bar{p}}$) production, using the predictions (16). Since the multiplicity associated with protons and antiprotons should be the same in the absence of stopping, we get

$$R = \frac{n^p}{n^{\bar{p}}} = (1 - f_{st}) + f_{st} \left(\frac{n^{(1)}}{n^{\bar{p}}} \right), \quad (20)$$

where we have omitted the contribution of the double stopping process, which is asymptotically suppressed at high energies according to (10), but may still be important at ISR energies [8]. Assuming that the multiplicity associated with antiprotons does not differ substantially from the average multiplicity of an inelastic event, $n^{\bar{p}} \simeq n^{inel}$, in accord with (13) and with experimental data [3], we obtain from (20) and (16) an estimate

$$R \simeq 1.05. \quad (21)$$

Even though the value (21) agrees within experimental errors with the measured value of $R \simeq 1.1$ [3], we may conclude that the contribution of double baryon stopping with higher associated multiplicity (16) is possible. A detailed experimental study of double baryon production in pp collisions would therefore be useful to clarify the situation.

Extrapolation to the pp collisions at the energies of $\sqrt{s} \simeq 6 \text{ TeV}$ (corresponding to the c.m.s. energy per nucleon-nucleon collision in Pb-Pb interactions at LHC) according to formulae (11, 13) with $\Delta = 0.08$ [12] yields then the following stopping fraction:

$$f_{st}(\sqrt{s} = 6 \text{ TeV}) \sim 5\%. \quad (22)$$

The nucleus-nucleus collisions will be accompanied by much larger stopping, as we discussed above, and we expect that the estimate (22) in this case can only be considered as a lower bound on the baryon asymmetry. Therefore we may expect substantial excess of baryons over antibaryons in the central rapidity region of nucleus-nucleus collisions at LHC.

It is interesting that already at SPS energy, the otherwise successful phenomenological approaches based on the topological expansion [13-16], but not taking into account the presence of the string junction explicitly, seem to underestimate [17] the pronounced baryon stopping observed experimentally in nucleus-nucleus collisions [18].

It would be useful to analyse the dynamics of baryon stopping in an approach where the topological structure of the baryon is explicit: the Skyrme model. The formation of baryon-antibaryon pairs in this approach is treated as the formation of topological defects in the quark condensate [19] (the net baryon number of the produced pairs is of course equal to zero). The picture proposed above

would invoke into consideration also the high-energy scattering of such topological defects. Since the Lorentz boost of such configurations is not trivial, one may expect the occurrence of the final states in which one or two Skyrmions are stopped in the central region, and the part of their pion field is “shaken off” to the fragmentation region. Such events would contribute to the baryon asymmetry in the central region. We leave the consideration of the baryon stopping dynamics in topological models for further studies.

Amazingly, the so-called “Centauro” and “Chiron” events reported in cosmic-ray emulsion experiments [20] and interpreted in favour of the existence of dense quark matter by Bjorken and McLerran [21], are characterized by non-vanishing baryon number density. As was formulated in ref. [22], “the fluid in the central region has no net baryon number, so that there would need to be a spontaneous generation of net baryon density to make these objects”. Indeed, in the scenario proposed by Bjorken [22], at ultra-relativistic energies the valence quarks of the colliding nuclei pass through each other, leaving behind a “little Universe” of zero net baryon density, which at the moment of its production contains a gluon sea. In our picture, this sea from the very beginning is made stormy by the presence of non-perturbative twists - specific configurations of the gluon field, which trace non-zero *net* baryon number. The “little Universe”, just like our big one, may therefore generate substantial baryon asymmetry.

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References

- [1] G.C. Rossi and G. Veneziano, Nucl. Phys. **B123** (1977) 507; Phys. Rep. **63** (1980) 153.
- [2] G. Veneziano, Nucl.Phys. **B74** (1974) 365; Phys.Lett. **B52** (1974) 220.
- [3] G. Belletini et al., Nuovo Cimento **42A** (1977) 85.
- [4] B. Alper et al., Nucl. Phys. **B100** (1975) 237.
- [5] G. Bendiscioli and D. Kharzeev, Riv. Nuovo Cimento **17** (1994) No.6.
- [6] A.H. Mueller, Phys. Rev. **D2** (1970) 2963.
- [7] O.V. Kancheli, JETP Lett. **11** (1970) 397.
- [8] L. Camilleri, Phys. Rep. **144** (1987) 51.
- [9] B.Z. Kopeliovich and B.G. Zakharov, Z. Phys. **C43** (1989) 241.
- [10] B. Alper et al., Phys. Lett. **B47** (1973) 275.
- [11] A.M. Rossi et al., Nucl. Phys. **B84** (1975) 269.
- [12] A. Donnachie and P.V. Landshoff, Phys. Lett. **B296** (1992) 227.
- [13] A. Capella and J. Trân Thành Van, Phys. Lett. **B114** (1982) 450.
- [14] A.B. Kaidalov, Phys. Lett. **B116** (1982) 459;
A.B. Kaidalov and K.A. Ter-Martirosyan, Sov.J.Nucl.Phys. **39** (1984) 1545.
- [15] G. Cohen-Tannoudji, A.E. Hassouni, J. Kalinowski and R. Peschanski, Phys. Rev. **D19** (1979) 3397.
- [16] A. Capella, U. Sukhatme, C.-I. Tan and J. Trân Thành Van, Phys. Rep. **236** (1994) 225.
- [17] A. Capella, private communication.
- [18] See, for example,
The NA35 Collaboration, M. Gaździcki et al., Nucl. Phys. **A590** 197c;
The NA49 Collaboration, S. Margetis et al., Nucl. Phys. **A590** 355c.
- [19] J. Ellis and H. Kowalski, Phys. Lett. **B214** (1988) 161; Nucl. Phys. **B327** (1989) 32;
J. Ellis, U. Heinz and H. Kowalski, Phys. Lett. **B214** (1988) 161;
J.I. Kapusta and A.M. Srivastava, Phys.Rev.**D52** (1995) 2977;
see also T.A. DeGrand, Phys.Rev.**D30** (1984) 2001.

- [20] Brazil-Japan Emulsion Chamber Collaboration, *unpublished*;
C.M.G. Lattes, Y. Fujimoto and S. Hasegawa, Phys.Rep. **65** (1980) 151.
- [21] J. Bjorken and L. McLerran, Phys.Rev.**D20** (1979) 2353.
- [22] J.D. Bjorken, Phys.Rev. **D27** (1983) 140.

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